

A composite light scalar, electro-weak symmetry breaking and the recent LHC searches.

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Abstract

We construct a model in which electro-weak symmetry breaking is induced by a strongly coupled sector, which is described in terms of a five-dimensional model in the spirit of the bottom-up approach to holography. We compute the precision electro-weak parameters, and identify regions of parameter space allowed by indirect tests. We compute the spectrum of scalar and vector resonances, which contains a set of parametrically light states that can be identified with the electroweak gauge bosons and a light dilaton. There is then a little desert, up to 2-3 TeV, where towers of resonances of the vector, axial-vector and scalar particles appear.

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1 Introduction.

The biggest open problem in high energy physics is to understand the origin of electro-weak symmetry breaking (EWSB) and hence of mass generation in the Standard Model (SM). Some new physics (particles and interactions) must be present at the electro-weak scale. The main physics goal of the Large Hadron Collider (LHC) is to collect as much experimental information as possible about electro-weak scale new physics, in order to address this problem.

From a theoretical point of view, the most fundamental question is whether the new interactions are weakly coupled, and hence EWSB originates from the condensation of elementary scalar fields (as is the case in the minimal version of the SM), or whether the new interaction is strong enough to trigger the formation of non-trivial composite condensates, as is the case in QCD regarding chiral symmetry breaking (as in Technicolor and its variations [1, 2, 3, 4]). From the experimental point of view, the most urgent question is whether a light scalar particle is present in the spectrum, the couplings of which are related to the masses of the other particles. In the (weakly coupled) minimal version of the SM, this is the Higgs particle. In certain (strongly coupled) models of dynamical electro-weak symmetry breaking, a particle with similar properties may be present as a consequence of the spontaneous breaking of (approximate) scale invariance: the dilaton [5, 6, 7, 8].

The main purpose of this paper is to provide a simple example of a model of the second class: there is no elementary scalar field, the fundamental interactions are strongly-coupled, and EWSB is due to the formation of some non-trivial condensate. However, because the model is approximately scale invariant, the condensate breaks spontaneously such symmetry, and as a result the spectrum contains a light dilaton.

Before the LHC program started, indirect searches for new physics have been carried out for quite some time, after the discovery of the W and Z gauge bosons. The surprising result of these searches is that, in spite of the fact that new interactions and new particles must be responsible for the weak scale $v \simeq 246$ GeV, none of these particles produces a sizable contribution to precision physics observables [9, 10], with the exception of a, possibly light, Higgs particle. This suggests the existence of a little desert in the spectrum, with a light Higgs followed by a gap, and possible new particles appearing only with masses in the 2–5 TeV

range. The dynamical generation and stabilization of such a gap is often referred to as the *little hierarchy* problem.

In order for a strongly-coupled model to be viable, we must be able to compute the precision parameters, the mass of the light states such as the W , Z and the possible light scalars, but also the mass spectrum of the heavy resonances, and show that all of these are in agreement with experimental bounds. This is a quite challenging program.

Gauge/gravity dualities [11, 12, 13] allow, in principle, for such calculations to be done. The idea is that one constructs a consistent background of an extra-dimension theory of gravity (string theory), and use it to compute quantities, such as correlation functions, in the dual four-dimensional gauge theory without ever referring to the microscopic Lagrangian of the latter. A difficulty of this approach is that it is still hard to study the dual of confining gauge theories. Notable examples exist, based on the deformations of the sphere [14, 15], of the conifold (a sample of the vast literature on the subject includes for instance [16, 17, 18, 19, 20, 21, 22, 23]) and of the Witten-Sakai-Sugimoto system [24]. But the systematics of all of these models is difficult to understand, hard to generalize, and calculations are technically challenging.

A simpler approach to holography consists of the so called bottom-up phenomenological models. Rather than relying on a fully consistent stringy construction, these models assume a set of simplifying assumptions for a five-dimensional system, and then proceed to perform the same kind of calculations one would do in a consistent gravity background. The advantage is obvious: constructing such models is much easier, and one can easily study whole classes of models by just dialing the free parameters. The disadvantage is also clear: not being derived from a fundamental theory of gravity, these models are somewhat arbitrary, and contain many more free parameters. This is reminiscent of what happens in effective theories, where one deals with large number of free parameters, all of which have the same fundamental origin, and hence would be correlated, if one could compute them from first principles.

Nevertheless, these models, because of their simplicity, are a very powerful tool, because they allow to ask how general certain results are, without having to rely on specific points of a parameter space that is otherwise hard to explore. The literature on the topic is rather extensive (for examples related to technicolor see [25]). We are going to build one such a model. The main differences with respect to what done in other papers is that the background is not AdS, and that we use a construction which is simple enough to allow to study at the same time both the spin-1 and spin-0 sectors.

The five-dimensional gravity theory is coupled to a scalar field, for which we assume a peculiarly simple (super)potential. We solve the coupled equations and find that the metric is asymptotically AdS both in the far UV and in the deep-IR, resembling the dual description of the field theory flow between two fixed points (see for example [26, 27]). We will see that the results and the numerology of such a setup are indeed not dissimilar from the rigorous results related to the PW flow (for example, our spectrum resembles qualitatively the one obtained in [28]). Because confinement is the hard thing to model, we introduce a hard cutoff in the IR, which introduces a scale Λ_0 in the theory. With abuse of language, we will call this the confinement scale, although in a rigorous sense this is rather just an IR regulator.

We then proceed to introduce the electro-weak group and electro-weak symmetry breaking. In the top-down approach, this can be done by introducing probe-branes, and showing that the curvature yields a U -shape embedding for the brane which can be interpreted in terms of the spontaneous breaking $U(1) \times U(1) \rightarrow U(1)$ (along the lines of what proposed in [24]). We mimic this procedure, by allowing the gauge bosons of $U(1)_L \times U(1)_R$ to propagate in the bulk, and imposing different boundary conditions in the IR for the axial and vector combination of the two fields, hence inducing the desired breaking pattern. Notice that the fact of not gauging a $SU(2)_L$ means that there are no W bosons in the spectrum. However, it would be trivial to extend the construction, at the only price of complicating the algebra. For our purposes, this is going to be satisfactory.

We then compute the spectrum of the four-dimensional spin-1 states. After holographic renormalization [29], it consists of two light states (that are identified with the massless photon and the massive Z boson), and two towers of heavy composite spin-1 states, the techni- ρ (vectorial) and techni- a_1 (axial) mesons, with masses controlled by the scale Λ_0 . We also compute the oblique precision parameters, compare to the experimental bounds, and hence fix some parameter in the model.

We then move on to study the scalar fluctuations of the background. This non-trivial task is made peculiarly simple by the fact that the formalism for doing so has been developed in detail in the literature [8, 30, 31], and by the fact that in the presence of one scalar only, with dynamics controlled by a superpotential, the equations become manageable. We compute the spectrum of perturbations, and observe that it always consists of a tower of heavy states, with mass controlled by Λ_0 , supplemented by an additional, parametrically light state (the dilaton).

We discuss in some detail under what circumstances such a dilaton is light. Its presence is due to the fact that the IR cutoff effectively amounts to the spontaneous breaking of scale invariance. The fact that the background has a non-trivial profile, interpolating between two AdS spaces with different curvature, can be interpreted (depending on the parameters) in terms of the presence of explicit breaking of scale invariance, which induces a mass. However, notice that because of the rough way in which we model confinement, our result is to be understood as a lower bound: other sources of explicit breaking will be present in a complete confining theory. It is remarkable that light states have been seen within the top-down approach (for example in [32]), but as we said, the systematics of this is not well understood.

We conclude the paper with numerical examples. We show that it is remarkably easy to produce a model in which the precision parameters are within the experimental bounds, in which a light scalar has a mass of the order of the Z boson, and in which a little desert follows, with heavy resonances appearing with masses above a few TeV.

The main message of the paper is that discovering (or not) a light scalar, with couplings similar to the Higgs particle of the Standard Model, does not imply that EWSB is (or is not) due to a new weakly coupled interaction. We want to show that it is actually easy to construct strongly coupled models that mimic such a scenario. In order to distinguish the two scenarios, the LHC experiments should explore in detail the whole mass region up to 2-4 TeV range, at high luminosity, and see whether other new resonances are there. A task within reach, provided the LHC increases its operational energy and luminosity, all of which is according to the future plan of the LHC program.

2 General set-up.

The model we study consists of a five dimensional theory, in which gravity, a scalar field Φ , and the $U(1) \times U(1)$ gauge bosons (L_M and R_M) propagate in a background geometry. We use the following conventions, as in [31], and follow the same formalism. Capital roman indexes $M = 0, 1, 2, 3, 5$ are five-dimensional space-time indexes, while greek indexes $\mu = 0, 1, 2, 3$ are restricted to the 4-dimensional Minkowski slices of the space. In this way, we label the space-time coordinates as $x^M = (x^\mu, r)$, with r the radial (fifth) direction. We write the five-dimensional metric g_{MN} with signature $\{- + + + +\}$. We assume that the background metric satisfies the ansatz

$$ds^2 = e^{2A} \eta_{\mu\nu} dx^\mu dx^\nu + dr^2, \quad (1)$$

with $A = A(r)$. Finally, we assume that the radial direction is compact, with $r_1 < r < r_2$. The boundaries are interpreted as an IR cutoff r_1 (which models in a very crude way the confinement scale of the putative four-dimensional dual field theory), and a UV cutoff r_2 (which is interpreted as a regulator, and will be removed at the end of the calculations via holographic renormalization).

The geometry is determined by the dynamics of the scalar field Φ , which represents a simple example of five-dimensional sigma-model coupled to gravity. We write the action as

$$\begin{aligned} \mathcal{S} \equiv & \int d^4x dr \left\{ \sqrt{-g} \Theta \left[\frac{1}{4} R + \mathcal{L}_5(\Phi, \partial_M \Phi, g_{MN}) \right] \right. \\ & + \sqrt{-\tilde{g}} \delta(r - r_1) [c_K K + \mathcal{L}_1(\Phi, \partial_\mu \Phi, \tilde{g}_{MN})] \\ & \left. - \sqrt{-\tilde{g}} \delta(r - r_2) [c_K K + \mathcal{L}_2(\Phi, \partial_\mu \Phi, \tilde{g}_{MN})] \right\}, \end{aligned} \quad (2)$$

where R is the Ricci scalar, where K is the extrinsic curvature, where the coupling $c_K = -1/2$ is fixed by consistency and where \mathcal{L}_i are the sigma-model actions. The step function is defined by $\Theta \equiv \Theta(r - r_1) - \Theta(r - r_2)$.

We define the action of the matter fields in terms of the real scalar field $\Phi = \Phi(x^\mu, r)$ as

$$\mathcal{L}_5 \equiv -\frac{1}{2} g^{MN} \partial_M \Phi \partial_N \Phi - V(\Phi), \quad (3)$$

$$\mathcal{L}_1 \equiv -\lambda_{(1)}(\Phi), \quad (4)$$

$$\mathcal{L}_2 \equiv -\lambda_{(2)}(\Phi). \quad (5)$$

For simplicity, we take the boundary actions for the scalar to depend only on the field, and not on the derivatives.

Assuming that there exists a superpotential W such that the potential V can be rewritten as

$$V = \frac{1}{2} (\partial_\Phi W)^2 - \frac{4}{3} W^2, \quad (6)$$

where $\partial_\Phi W \equiv \frac{\partial W}{\partial \Phi}$, then the system can be reduced to first-order equations

$$A' = -\frac{2}{3} W, \quad (7)$$

$$\bar{\Phi}' = \partial_\Phi W, \quad (8)$$

in the sense that all solutions to Eqs. (7)-(8) are also solutions to the second-order equations derived from \mathcal{S} .¹

We call $\Phi_1 = \Phi(r_1)$ and $\Phi_2 = \Phi(r_2)$ the boundary values of the scalar, which we use to define the boundary conditions. Accordingly, the general form of the localized potentials must be

$$\lambda_{(1)} = W(\Phi_1) + \partial_\Phi W(\Phi_1)(\Phi - \Phi_1) + \frac{1}{2} \partial_\Phi^2 \lambda_{(1)}(\Phi - \Phi_1)^2, \quad (9)$$

$$\lambda_{(2)} = W(\Phi_2) + \partial_\Phi W(\Phi_2)(\Phi - \Phi_2) + \frac{1}{2} \partial_\Phi^2 \lambda_{(2)}(\Phi - \Phi_2)^2. \quad (10)$$

¹Given a five-dimensional gauged supergravity for which a superpotential W is known, the first-order equations Eqs. (7)-(8) agree with the BPS equations one would obtain by studying the supersymmetry variations of the gravitino and dilatino. Hence, solving Eqs. (7)-(8) would amount to finding supersymmetric backgrounds. The converse is not true: the existence of W , and of solution to Eqs. (7)-(8), does not imply that the model is supersymmetric. In the present context, the fact that we study a model defined by a given choice of W is to be understood as just a technical tool, yielding a simple solution-generating technique. Notice in particular that the model we study does not contain any fermions, and that the fluctuations of the scalars are described by second-order equations. Our study does not assume nor imply that the strongly-coupled dual field theory be supersymmetric, in particular the existence of a light composite scalar is not due to supersymmetry.

In this way, the whole system can be described in terms of first order equations which satisfy automatically the second-order equations derived by variational principle from the five-dimensional action.

For reasons that will become clear later, we find it convenient to define:

$$N \equiv \partial_\Phi^2 W - \frac{(\partial_\Phi W)^2}{W}, \quad (11)$$

$$\partial_\Phi^2 \lambda_{(1)} \equiv \partial_\Phi^2 W|_{r_1} + (m_1^2), \quad (12)$$

$$\partial_\Phi^2 \lambda_{(2)} \equiv \partial_\Phi^2 W|_{r_2} - (m_2^2), \quad (13)$$

where m_i^2 are arbitrary mass terms that we will (conservatively) choose to diverge $m_i^2 \rightarrow +\infty$.

The gauge bosons are introduced as probes, which are not back-reacted on the metric, via the following action:

$$\mathcal{S}_5 \propto -\frac{1}{4} \int d^4x \int_{r_1}^{r_2} dr \left[\left(a(r) + b(r) D \delta(r - r_2) \right) \eta^{\mu\nu} \eta^{\gamma\delta} F_{\mu\gamma} F_{\nu\delta} + 2b(r) \eta^{\mu\nu} F_{r\mu} F_{r\nu} \right], \quad (14)$$

where F stands for the field-strength of both the V and A fields, which are defined by

$$L = \cos \theta_W A + \sin \theta_W V, \quad (15)$$

$$R = -\sin \theta_W A + \cos \theta_W V. \quad (16)$$

The gauge couplings g and g' of the $U(1)_L$ and $U(1)_R$, respectively, determine the mixing angle $\tan \theta_W = \frac{g'}{g}$ that relates the (L, R) basis and the (V, A) basis.

After Fourier-transforming in the four Minkowski directions (and writing the four-momentum q^μ with the opposite convention with signature $(+, -, -, -)$ for later convenience), we can rewrite $A^\mu(q, r) \rightarrow A^\mu(q) v(q, r)$, in such a way that all the information about the bulk theory be contained in the functions $v(q, r)$. The two kinds of fields have the same equations of motion, given by

$$\partial_r \left(b(r) \partial_r v(q, r) \right) = -q^2 a(r) v(q, r). \quad (17)$$

Electro-weak symmetry breaking is induced by the fact that the solutions $v_{A,V}$ must satisfy *different* IR boundary conditions:

$$v_A(r_1) = 0 = \partial_r v_V(r_1). \quad (18)$$

For the functions a and b we take the simplest possible form, which can be thought of as coming from allowing the spin-1 fields to propagate in the space with the geometry determined by the five-dimensional sigma-model background:

$$a(\rho) = \sqrt{-g} g^{xx} g^{xx} = 1, \quad (19)$$

$$b(\rho) = \sqrt{-g} g^{xx} g^{rr} = e^{2A}. \quad (20)$$

3 The model.

We borrow model C from [31], for which

$$W = -\frac{3}{2} - \frac{\Delta}{2} \Phi^2 + \frac{\Delta}{3\Phi_I} \Phi^3, \quad (21)$$

and we focus on the solution

$$\bar{\Phi} = \frac{\Phi_I}{e^{\Delta(r-r_*)} + 1}, \quad (22)$$

which as expected is a 1-parameter family labelled by r_* .

In this case, evaluating on the solution, one has

$$W = \frac{1}{6} \left(-9 - \frac{\Delta(1 + 3e^{\Delta(r-r_*)})\Phi_I^2}{(1 + e^{\Delta(r-r_*)})^3} \right), \quad (23)$$

$$W' = -\frac{\Delta e^{\Delta(r+r_*)}\Phi_I}{(e^{\Delta r} + e^{\Delta r_*})^2}, \quad (24)$$

$$A = \frac{1}{9} \left(9r + \Phi_I^2 \frac{e^{\Delta(r+r_*)}}{(e^{\Delta r} + e^{\Delta r_*})^2} + \Phi_I^2 \Delta r - \Phi_I^2 \log[1 + e^{\Delta(r-r_*)}] \right) - \frac{1}{9} \Phi_I^2 \left(\frac{1}{2 \cosh(\Delta r_*) + 2} - \log(1 + e^{-\Delta r_*}) \right). \quad (25)$$

Notice that we conventionally fix $r_1 = 0$ from now on, and we chose an integration constant in A so that $A(0) = 0$. In this way, all the masses will be expressed in units of an overall scale Λ_0 , that will be fixed by the mass of the Z boson.

Before we start to study the physics of the model we do some parameter counting. The parameters in the model are: r_1 , r_2 , Φ_I , Δ , $m_{1,2}^2$, r_* , D and an overall constant in front of \mathcal{S}_5 . The constant in front of \mathcal{S}_5 is related to the effective five-dimensional gauge coupling, and we will later call it \mathcal{N} . This constant together with D effectively control the strength of the gauge couplings of the light degrees of freedom (photon and Z boson), and of the resonances. Because we take the gauge couplings of the lightest states to be identified with the standard-model couplings, there is only one free parameter, which is the ratio of these gauge couplings to the effective couplings of the resonances (one can think of the latter as the equivalent to the $g_{\rho\pi\pi}$ coupling of QCD). We set $r_1 = 0$, hence setting all mass scales in terms of the same scale, which enters both into quantities related to confinement and to symmetry breaking. r_2 is a UV-regulator, which will be taken to infinity at the end of the calculations, and hence has no physical meaning. The boundary masses $m_{1,2}^2$ do not enter in the equations determining the background, but only in the equations for the scalar fluctuations. We will take them to diverge, as anticipated, as a conservative assumption (if we find a light state in the spectrum of scalar perturbations in this case, it will in general be parametrically light also for any other choice of $m_{1,2}^2$ [31]). Concluding: there are four parameters that are physical and free: Φ_I , which controls the departure from AdS of the bulk geometry, Δ which can be interpreted in terms of the dimensionality of the dual operator inducing the departure from conformality, r_* which controls the physical scale at which the departure from conformality arises, and which is independent of the confinement scale, and the effective coupling of the spin-1 resonances, which can be interpreted in terms of the large- N expansion parameter of the dual field theory.

A comment about the model. The remarkable simplicity of the superpotential clearly indicates that this is not a model that has been obtained from a complete string theory with the familiar tools of consistent truncation. However, the background that it yields has many features that are very similar to those obtained in more rigorous constructions (for example, see [26]). The simplicity of this model allows us to perform the actual calculations (almost) analytically, and we expect that the main features we obtain could be derived from a large class of string-motivated models.

4 Electroweak symmetry breaking.

What we are constructing is the low energy description of a simplified model of dynamical electroweak symmetry breaking, in which as anticipated we simplify the standard model by replacing the $SU(2)_L \times U(1)_Y \rightarrow U(1)_{\text{e.m.}}$ with $U(1)_L \times U(1)_R \rightarrow U(1)_V$, with gauge couplings g and g' for the L and R gauge groups, respectively. The Lagrangian of the two light gauge bosons, focusing on the transverse vacuum polarizations, can be written in general as

$$\mathcal{L}_4 = -\frac{1}{2}P^{\mu\nu}A_\mu^i(q)A_\nu^j(-q)\pi(q^2)_{ij}, \quad (26)$$

where $i = L, R$, where $P^{\mu\nu} = \eta^{\mu\nu} - q^\mu q^\nu / q^2$ and where the $\pi^{i,j}$ are (possibly complicated) functions of the four-momentum. Such functions are going to be non-analytical in general. However, under reasonable assumptions, they will be regular at the origin, and can be expanded as

$$\pi(q^2) = \pi(0) + q^2\pi'(0) + \frac{1}{2}(q^2)^2\pi''(0) + \dots. \quad (27)$$

The precision electro-weak parameters are defined by the following relations [10]:

$$\hat{S} = \frac{g}{g'}\pi'_{LR}(0), \quad (28)$$

$$X = \frac{1}{2}M_W^2\pi''_{LR}(0), \quad (29)$$

$$W = \frac{1}{2}M_W^2\pi''_{LL}(0), \quad (30)$$

$$Y = \frac{1}{2}M_W^2\pi''_{RR}(0), \quad (31)$$

where M_Z is the mass of the lightest axial state, while we just define $M_W^2 \equiv \cos^2\theta_W M_Z^2$. Notice that because we restrict our attention to the neutral gauge bosons, the custodial-symmetry breaking parameters \hat{T} and \hat{U} do not exist.² In these definitions, one assumes that the fields are normalized so that $\pi'_{LL}(0) = 1 = \pi'_{RR}(0)$. By making use of the rotations from the (L, R) basis to the vector axial-vector basis (V, A) , we can rewrite

$$\hat{S} = \cos^2\theta_W \left(\pi'_V(0) - \pi'_A(0) \right), \quad (32)$$

$$X = \cos^2\theta_W \frac{M_Z^2}{2} \sin\theta_W \cos\theta_W \left(\pi''_V(0) - \pi''_A(0) \right), \quad (33)$$

$$W = \cos^2\theta_W \frac{M_Z^2}{2} \left(\sin^2\theta_W \pi''_V(0) + \cos^2\theta_W \pi''_A(0) \right), \quad (34)$$

$$Y = \cos^2\theta_W \frac{M_Z^2}{2} \left(\cos^2\theta_W \pi''_V(0) + \sin^2\theta_W \pi''_A(0) \right). \quad (35)$$

All of this means that having the functions $\pi_{V,A}$ up to the second derivative in q^2 is enough to yield all the precision parameters.

²In order to avoid generating a significant contribution to electroweak precision parameters related to custodial symmetry breaking, one might have to extend the bulk symmetry to $SU(2) \times SU(2)$. In doing so one would ensure that the \hat{T} and \hat{U} parameters are negligibly small, independently of the bulk dynamics. See for example [33], and references therein, for a general discussion of this and related issues, which we do not address in our study.

The fit of the experimental data from precision electroweak physics indicates that all of these four precision parameters must be at most of the order of 3×10^{-3} [10]. The precise results depend on what confidence level we use and on what value of the reference mass for the Higgs particle we adopt, and the various bounds are correlated. For the present purposes, this very rough bound is sufficient. One can see that (at least approximately), provided the precision parameters are all small (which in turns means that $\pi'_A \simeq \pi'_V = 1$) the symmetry-breaking scale can be measured by the mass of the lightest axial mode:

$$M_Z^2 \simeq -\frac{\pi_A(0)}{\pi'_A(0)} \simeq -\pi_A(0). \quad (36)$$

Before focusing on the specific model of this paper, let us remind the reader about how to perform the calculation of the precision parameters in the general case of a five-dimensional model of this type. For the vectorial bosons, it is convenient to define the following:

$$\partial_r v(q, r) \equiv \gamma(q, r) v(q, r), \quad (37)$$

and replace in the bulk equations, which yield

$$\partial_r(b\gamma) + b\gamma^2 + q^2 a = 0, \quad (38)$$

where we simplified the notation by dropping obvious terms. This first-order equation contains the same amount of information as the original second-order equation, but it has the advantage that the overall normalization of the state v drops. The functions γ must satisfy the boundary condition:

$$\gamma(q, r_1) = 0. \quad (39)$$

For the axial part, it is convenient to define the inverse of the function γ , via

$$v(q, r) \equiv X(q, r) \partial_r v(q, r), \quad (40)$$

which obeys the (same) equation

$$-b^2 \partial_r \left(\frac{X}{b} \right) + b + q^2 a X^2 = 0, \quad (41)$$

subject to the (different) boundary condition

$$X(q, r_1) = 0. \quad (42)$$

One hence must find solutions to the equations for γ and X , subject to the boundary conditions in the deep-IR. Finally, the physical spectrum and the two-point correlation-functions can be studied from evaluating the solutions at the UV-boundary r_2 , and then trying to take the limit $r_2 \rightarrow +\infty$, if possible. Before doing so, some more about solving the actual equations.

For the calculation of the precision parameters, we need only the first few terms of the expansion of γ and X in q^2 . We define the low-energy expansions

$$\gamma(q, r) \equiv \gamma_0(\rho) + q^2 \gamma_1(\rho) + \frac{1}{2} (q^2)^2 \gamma_2(\rho) + \dots, \quad (43)$$

$$X(q, \rho) \equiv X_0(\rho) + q^2 X_1(\rho) + \frac{1}{2} (q^2)^2 X_2(\rho) + \dots, \quad (44)$$

and expand accordingly the equations of motion which yield

$$\partial_r(b\gamma_0) + b\gamma_0^2 = 0, \quad (45)$$

$$\partial_r(b\gamma_1) + 2b\gamma_0\gamma_1 + a = 0, \quad (46)$$

$$\partial_r(b\gamma_2) + 2b(\gamma_0\gamma_2 + \gamma_1^2) = 0, \quad (47)$$

$$\dots,$$

$$-b^2\partial_r\left(\frac{X_0}{b}\right) + b = 0, \quad (48)$$

$$-b^2\partial_r\left(\frac{X_1}{b}\right) + aX_0^2 = 0, \quad (49)$$

$$-b^2\partial_r\left(\frac{X_2}{b}\right) + 4aX_0X_1 = 0, \quad (50)$$

$$\dots,$$

subject to the boundary conditions

$$\gamma_i(r_1) = 0, \quad (51)$$

$$X_i(r_1) = 0, \quad (52)$$

which ensure that the IR boundary conditions are satisfied for every q . A further simplification occurs thanks to the boundary conditions, which can be satisfied only provided $\gamma_0 = 0$ identically for any r . This reflects the physical fact that γ is going to be related to the inverse of the propagator of the gauge boson of an unbroken $U(1)$, and hence must contain a massless mode.

At least formally, the equations can be integrated, to yield:

$$\gamma_0(r) = 0, \quad (53)$$

$$\gamma_1(r) = -\frac{1}{b(r)} \int_{r_1}^r d\omega a(\omega), \quad (54)$$

$$\gamma_2(r) = -\frac{1}{b(r)} \int_{r_1}^r d\omega 2b(\omega)\gamma_1(\omega)^2, \quad (55)$$

$$\dots$$

$$X_0(r) = b(\rho) \int_{r_1}^r \frac{d\omega}{b(\omega)}, \quad (56)$$

$$X_1(r) = b(\rho) \int_{r_1}^r \frac{a(\omega)X_0(\omega)^2}{b(\omega)^2}, \quad (57)$$

$$X_2(r) = b(\rho) \int_{r_1}^r \frac{4a(\omega)X_0(\omega)X_1(\omega)}{b(\omega)^2}, \quad (58)$$

$$\dots$$

We can now evaluate the effective five-dimensional action on-shell. Doing so, because we are satisfying the bulk equations and the IR boundary conditions, leaves only a boundary term localized at the UV boundary r_2 . To this, we must add localized boundary actions, which depend only on the four-dimensional

fields A_μ and V_μ . The results are

$$\pi_V(0) = 0, \quad (59)$$

$$\pi_A(0) = \mathcal{N}(r_2) \left(C(r_2)b(r_2) + b(r_2)\frac{1}{X_0(r_2)} \right), \quad (60)$$

$$\pi'_V(0) = \mathcal{N}(r_2) \left(D(r_2)b(r_2) + b(r_2)\gamma_1(r_2) \right), \quad (61)$$

$$\pi'_A(0) = \mathcal{N}(r_2) \left(D(r_2)b(r_2) - b(r_2)\frac{X_1(r_2)}{X_0(r_2)^2} \right), \quad (62)$$

$$\pi''_V(0) = \mathcal{N}(r_2) \left(E(r_2)b(r_2) + b(r_2)\gamma_2(r_2) \right), \quad (63)$$

$$\pi''_A(0) = \mathcal{N}(r_2) \left(E(r_2)b(r_2) + b(r_2) \left(\frac{2X_1(r_2)^2}{X_0(r_2)^3} - \frac{X_2(r_2)}{X_0(r_2)^2} \right) \right), \quad (64)$$

where \mathcal{N} is a normalization that is related to the wave-function normalization of the four-dimensional fields, and that is fixed by requiring that the photon be canonically normalized, while C , D and E are counterterms. Notice however that the counterterm C naively breaks gauge-invariance, and hence must come (if non-vanishing) from a localized (elementary) Higgs field. The counterterm D is always present, and effectively controls the strength of the (weak) gauging of the global $U(1)$ s. The counterterm E comes from higher-order operators of the form $(\partial F)^2$, which would signal the fact that the theory is UV-incomplete.

Let us now move back to our example. In our case $C = E = 0$, $r_1 = 0$, $a = 1$ and $b = e^{2A}$. With all of this in place, one finds that

$$b(r_2)\gamma_1(r_2) = - \int_0^r d\omega = -r_2, \quad (65)$$

which signals the presence of a (logarithmic) divergence in the kinetic term of the vector fields. This is expected in full generality, and is connected with the running of the (weak) four-dimensional couplings because the elementary fields of the strongly-coupled dual theory are charged. This allows us to choose the parameter D to be

$$Db(r_2) = r_2 - \frac{1}{\varepsilon^2}, \quad (66)$$

effectively trading the divergence for a free parameter. With the additional choice of normalization

$$\mathcal{N} = -\varepsilon^2, \quad (67)$$

one recovers

$$\pi'_V = 1. \quad (68)$$

Hence, for generic values of the parameters of the model, we have

$$\pi_V(0) = 0, \quad (69)$$

$$\pi_A(0) = -\varepsilon^2 b(r_2) \left(\frac{1}{X_0(r_2)} \right) = \frac{-\varepsilon^2}{\int_0^{r_2} e^{-2A(\omega)} d\omega}, \quad (70)$$

$$\pi'_V(0) = 1, \quad (71)$$

$$\pi'_A(0) = -\varepsilon^2 \left(r_2 - \frac{1}{\varepsilon^2} - \frac{b(r_2)X_1(r_2)}{X_0(r_2)^2} \right), \quad (72)$$

$$\pi''_V(0) = -\varepsilon^2 b(r_2) \left(\gamma_2(r_2) \right), \quad (73)$$

$$\pi''_A(0) = -\varepsilon^2 b(r_2) \left(\frac{2X_1(r_2)^2}{X_0(r_2)^3} - \frac{X_2(r_2)}{X_0(r_2)^2} \right). \quad (74)$$

Notice in particular that because A is approximately linear at large- r , the integrals now all converge, and the limit $r_2 \rightarrow +\infty$ can be taken.

Let us do all the calculations explicitly. This can be done numerically, or with the following approximations put in place. It is reasonable (though not necessary) to assume that Δ lies in the range $1 < \Delta < 2$. The holographic interpretation of this is that the bulk is not AdS because the dual theory is deformed by the insertion of an operator of dimension $4 - \Delta$ (for a concrete example, see again [26]). It is also reasonable to assume that Φ_I be at most $\mathcal{O}(1)$. If so, the warp factor A can be approximated very well by

$$A \simeq r \left(\frac{\Delta \Phi_I^2}{9} + 1 \right) \Theta(r_* - r) \quad (75)$$

$$\begin{aligned} & + \Theta(r - r_*) \left(\frac{1}{9} \Delta r_* \Phi_I^2 - \frac{1}{9} \left(\frac{1}{2 \cosh(\Delta r_*) + 2} - \log(1 + e^{-\Delta r_*}) \right) \Phi_I^2 + r \right), \\ & \equiv w_1 r \Theta(r_* - r) + \Theta(r - r_*) (r + w_2), \end{aligned} \quad (76)$$

where the constants w_1 and w_2 have a very different meaning (and dimensionality). Notice that this is a very rough approximation, especially for small Δ and for $r \simeq r_*$.

Then, one finds that

$$\begin{aligned} X_0(r) &= e^{2A(r)} \int_0^r d\omega e^{-2A(\omega)} \\ &\simeq \frac{e^{2A(r)}}{2} \left[e^{-2(r_* + w_2)} - e^{-2(r + w_2)} + \frac{1 - e^{-2w_1 r_*}}{w_1} \right], \end{aligned} \quad (77)$$

provided $r > r_* > 0$, while

$$X_0(r) = e^{2A(r)} \int_0^r d\omega e^{-2A(\omega)} \simeq \frac{e^{2r}}{2} (1 - e^{-2r}) \quad (78)$$

when $r_* \ll 0$.

From these, by taking $r_2 \rightarrow \infty$, we can read

$$M_Z^2 \simeq -\pi_A(0) \simeq \frac{2\varepsilon^2 w_1}{1 - e^{-2w_1 r_*} + w_1 e^{-2(r_* + w_2)}}, \quad (79)$$

$$\simeq 2\varepsilon^2 w_1, \quad (80)$$

for $r_* \gg 0$ and

$$M_Z^2 \simeq 2\varepsilon^2, \quad (81)$$

when $r_* \ll 0$. This is also the limit $\Phi_I \rightarrow 0$, in which the space becomes AdS with unit curvature, and $M_Z^2 \simeq 2\varepsilon^2$.

The physical value of the Z mass is $M_Z^{exp} \simeq 91$ GeV. Equating the two, we find that $M_Z^{exp} = M_Z \Lambda_0$, where Λ_0 is the confinement scale controlled by r_1 . This allows us to conclude that for large r_* we have

$$\Lambda_0 = \frac{M_Z^{exp}}{\sqrt{2}w_1\varepsilon}, \quad (82)$$

and that in order to recover proper units we must multiply all the masses we find by this common scale.

Next, we compute

$$X_1(r) = e^{2A(r)} \int_0^r d\omega e^{-4A(\omega)} X_0(\omega)^2, \quad (83)$$

and from it find that the precision parameter \hat{S} , which provided $r_* > 0$ is given by

$$\begin{aligned} \hat{S} = & \frac{\cos^2 \theta_W \varepsilon^2}{4w_1 (e^{2r_* w_1} w_1 - e^{2(r_* + w_2)} + e^{2(w_1 r_* + r_* + w_2)})^2} \left(e^{4r_* w_1} (4r_* + 3) w_1^3 - 4e^{2(w_1 r_* + r_* + w_2)} (2r_* + 1) w_1^2 \right. \\ & \left. + 4e^{2(2w_1 r_* + r_* + w_2)} (2r_* + 1) w_1^2 + 3e^{4(w_1 r_* + r_* + w_2)} - 4e^{2r_* (w_1 + 2) + 4w_2} (2r_* w_1 + 1) + e^{4(r_* + w_2)} (4r_* w_1 + 1) \right), \end{aligned} \quad (84)$$

which for large values of r_* becomes

$$\hat{S} \simeq \frac{3\varepsilon^2 \cos^2 \theta_W}{4w_1}. \quad (85)$$

On the other hand, for $r_* \ll 0$ one recovers the pure AdS case with unit curvature, which yields

$$\hat{S} \simeq \frac{3\varepsilon^2 \cos^2 \theta_W}{4}. \quad (86)$$

At this point, we observe that for any reasonable values of $w_{1,2}$ and r_* , this differs very little from the pure-AdS case. Only when one chooses $\Phi_I \gg 1$ is the result suppressed by $\mathcal{O}(1)$ factors in respect to the pure-AdS case, but then many of the current approximations cannot be trusted.

This means that one has to choose small values of ε^2 in order for the bound $\hat{S} < 0.003$ to be satisfied, which in turns means that the parameters W and Y are automatically within their limit, being proportional to ε^4 . For example, if we choose $\varepsilon = 0.07$, $\Phi_I = 1$, $\Delta = 1$, and a somewhat largish value of r_* we obtain $\hat{S} \simeq 0.0026$, and $\Lambda_0 \simeq 873$ GeV.

5 Spectrum of heavy spin-1 states.

We want now to compute the spectrum of the heavy resonances in the model. In order to do so, we cannot rely on the formalism of the previous section, by making use of an expansion in q^2 , but instead we need to

reconsider the bulk equation for $\gamma(q, r)$. We must compute the zeros of the complete $\pi_{V,A}$:

$$\pi_V(q^2) = \lim_{r_2 \rightarrow +\infty} -\varepsilon \left(r_2 q^2 - \frac{1}{\varepsilon^2} q^2 + e^{2A(r_2)} \gamma(q, r_2) \right), \quad (87)$$

$$\pi_A(q^2) = \lim_{r_2 \rightarrow +\infty} -\varepsilon \left(r_2 q^2 - \frac{1}{\varepsilon^2} q^2 + e^{2A(r_2)} \frac{1}{X(q, r_2)} \right). \quad (88)$$

We do so approximately for $r_* \ll 0$, by assuming that $\Delta\Phi_I \leq 1$, in which case we can use the AdS result which yields:

$$\pi_V = q^2 + \varepsilon^2 q^2 \left(\frac{\pi Y_0(q)}{2J_0(q)} - \gamma_E - \log \frac{q}{2} \right), \quad (89)$$

$$\pi_A = q^2 + \varepsilon^2 q^2 \left(\frac{\pi Y_1(q)}{2J_1(q)} - \gamma_E - \log \frac{q}{2} \right), \quad (90)$$

where Y and J are Bessel functions, and γ_E is the Mascheroni constant. The opposite case in which r_* is very large, can be obtained by rescaling by a factor of w_1 all the dimensionful quantities, so that for parametrically large r_*

$$\pi_V = \frac{q^2}{w_1^2} + \varepsilon^2 \frac{q^2}{w_1^2} \left(\frac{\pi Y_0(q/w_1)}{2J_0(q/w_1)} - \gamma_E - \log \frac{q/w_1}{2} \right), \quad (91)$$

$$\pi_A = \frac{q^2}{w_1^2} + \varepsilon^2 \frac{q^2}{w_1^2} \left(\frac{\pi Y_1(q/w_1)}{2J_1(q/w_1)} - \gamma_E - \log \frac{q/w_1}{2} \right). \quad (92)$$

Under the assumption that r_* be large, the resulting spectrum consists of two towers of heavy states:

$$M_{Vn} \simeq \left(n - \frac{1}{4} \right) \pi w_1, \quad (93)$$

$$M_{An} \simeq \left(n + \frac{1}{4} \right) \pi w_1, \quad (94)$$

for $n = 1, 2, \dots$, plus a massless vectorial state $M_{V0} = 0$ (the photon) and a light axial vector state $M_{a0} = M_Z \simeq \sqrt{2}\varepsilon w_1$.

Notice that the w_1 -dependence does not drop even normalizing with the mass of the Z boson (as also clear from the fact that \hat{S} does depend on w_1 and in general on r_*). The first excited state of the vectorial mode (ρ meson) has a mass $M_{V1} \gtrsim 2.3$ TeV. For the first excited axial-vector mode we have $M_{a1} \gtrsim 3.6$ TeV. Notice that one could lower these masses by allowing for larger values of Φ_I , but in this case some of our approximations might fail. We are not going to do so in the following.

6 Scalar Spectrum.

We can also compute the spectrum of scalar perturbations of the sigma-model. We do so by using the formalism in [31], which requires to solve the equation

$$\left[e^{-4A} (\partial_r + N) e^{4A} (\partial_r - N) + e^{-2A} q^2 \right] \mathbf{a} = 0, \quad (95)$$

subject to the boundary conditions

$$\left[\frac{e^{2A}}{q^2} \frac{(\partial_\Phi W)^2}{W} \right] (\partial_r - N) \mathbf{a} \Big|_{r_i} = \mathbf{a} \Big|_{r_i} . \quad (96)$$

Notice that this comparatively simple results hold only because we took $m_i^2 \rightarrow +\infty$.

For the heavy states, and as long as the departure from pure AdS is small, we can approximately solve these equations by noticing that $N \simeq -\Delta$ in this limit, so that the heavy mass eigenstates can be found by solving

$$((\partial_r + 4 - \Delta)(\partial_r + \Delta) + e^{-2r} q^2) \mathbf{a} = 0 , \quad (97)$$

$$\mathbf{a} \Big|_{r_i} = 0 . \quad (98)$$

The general solution is

$$\mathbf{a} = e^{-2r} (b_1 J_{\Delta-2}(e^{-r} q) + b_2 Y_{\Delta-2}(e^{-r} q)) , \quad (99)$$

and the spectrum (for $r_1 = 0$ and $r_2 \rightarrow +\infty$) is approximately given by the zeros of $J_{\Delta-2}(q)$, which for $\Delta = 1$ are

$$M_n \simeq \left(n + \frac{1}{4} \right) \pi . \quad (100)$$

This is the same result as for the axial-vector case. Notice instead that for very large values of r_* the result is very different: the spectrum is controlled by the effective scaling dimension computed at the IR fixed point, which is very large, and hence the whole spectrum is shifted upwards. Furthermore, in this case one finds that there is an overall small rescaling, due to w_1 , in analogy to what happens for the heavy vectors.

More interesting is the case of the lightest state in the spectrum, which has to be treated more carefully. As known from the literature, for $\Delta > 2$, because we are going to take $r_2 \rightarrow +\infty$, this state would be exactly massless. More interesting is the case where $\Delta < 2$, which is what we focus on. For extreme values of $r_* \rightarrow \pm\infty$, the dilaton becomes massless. The complete r_* -dependence has been studied in details in [31] and agrees with the PW case [28]. It has a maximum, for r_* close to zero, and falls-off exponentially elsewhere. Because of this, it is actually possible in principle to use the mass of the light dilaton in order to extract information about the physical scale $\Lambda_* \equiv \Lambda_0 e^{r_*}$ that characterizes the background dynamics.

Let us conclude by numerically studying a few specific points in the parameter space of this model, as a way to both test our approximations, and to illustrate what is going on. We fix $\Delta = 1 = e^{r_1} = \Phi_I$ and $\varepsilon = .07$ on the basis of the approximations we made so far, and numerically compute all the relevant physical quantities. The results are summarized in Table 1 and in Fig. 1. The motivation for these choices is not particularly strong: there is no reason why Φ_I should be small, and hence we take it to be of the same order as the confinement scale $e^{r_1} = 1$. For the dimensionality, we choose it to be $\Delta = 1$. For example this is what happens in the Pilch-Warner flow [26], dual to the flow from $\mathcal{N} = 4$ SYM to the Leigh-Strassler conformal fixed point [27] (because of supersymmetry, in that case there are actually two active scalars in the five-dimensional theory, and their VEVs correspond to the deformation of the field theory by the presence of a mass terms for one of the fermions, with $\Delta = 1$, and one of the scalars, with $\Delta = 2$). Notice how the very simple model studied here produces a light scalar spectrum which is qualitatively very similar to what happens in the Pilch-Warner case [28], meaning that the bulk geometry we choose is realistic.

	$\varepsilon = 0.07, \Phi_I = 1, \Delta = 1$						
r_*	1.5	2.	2.2	2.3	2.4	2.5	2.6
M_Z	0.0912	0.0912	0.0912	0.0912	0.0912	0.0912	0.0912
m	0.2410	0.1731	0.1490	0.1378	0.1272	0.1172	0.1078
\hat{S}	0.0026	0.0026	0.0026	0.0026	0.0026	0.0026	0.0026
Λ_*	3.9	6.5	7.9	8.7	9.6	10.7	11.8
M_{V1}	2.3	2.3	2.3	2.3	2.3	2.3	2.3
M_{V2}	5.3	5.3	5.3	5.3	5.3	5.3	5.3
M_{V3}	8.2	8.3	8.3	8.3	8.3	8.3	8.3
M_{A1}	3.7	3.7	3.7	3.7	3.7	3.7	3.7
M_{A2}	6.7	6.7	6.7	6.7	6.8	6.8	6.8
M_{A3}	9.7	9.7	9.8	9.8	9.8	9.8	9.8
M_1	5.3	5.6	5.7	5.7	5.7	5.8	5.8
M_2	8.2	8.6	8.7	8.8	8.8	8.9	8.9

Table 1: Numerical results for $\Delta = 1 = \Phi_I = e^{r_1}$, $\varepsilon = 0.07$, $r_2 \rightarrow +\infty$ by varying r_* . All masses are in TeV.

Some brief comments. Because $\Delta < 2$, the mass of the light scalar state converges when $r_2 \rightarrow +\infty$ and stays finite [31]. However, it is suppressed as by r_* . The comparison we make with the current bounds on the Higgs mass has to be taken with a grain of salt. Our main point is that it is not particularly difficult to find regions of parameter space where the dilaton has a mass in the allowed low-mass range $114 < m < 141$ GeV [34], by dialing $r_* \sim 2.4$. It is also remarkable that there is an actual upper bound on $m \lesssim 350$ GeV. Because of the choice of ε , and hence because of the bound on \hat{S} , there is a little desert between m and the mass of the lightest ρ meson (around 2.3 TeV). Above this, one finds a rather complicated spectrum. As expected, the scalar resonances all show the effect of the scale r_* , while the spin-1 resonances are virtually insensitive to r_* .

7 Critical discussion.

Before concluding the paper, we want to alert the reader about two more intrinsic limitations of the procedure we followed, and suggest to use caution in interpreting the results.

The model we propose contains a light scalar in the spectrum. For illustration purposes, we compared its value, for a few choices of parameters, with the bounds on the mass of the standard-model Higgs particle. The reason why this is not completely unreasonable is that we know that the light state is a (pseudo-)dilaton, and hence its couplings are going to be very similar to those of the Higgs particle (which is a dilaton itself). Namely, the light scalar couples to the light SM particles with couplings proportional to their masses (for fermions and gauge bosons), or to the beta-functions computed by using the fields that are heavier than the dilaton itself (for the gluon-gluon and photon-photon interactions).

There is no reason to expect the overall magnitude of such couplings to be precisely the SM one: in general the physical particle results from a mixing between the fluctuation of the metric and of the background scalar field, and this will result in a suppression of the couplings (due to some mixing angle). However, the new strongly coupled sector contains (techni-)fermions that are electrically charged. This will

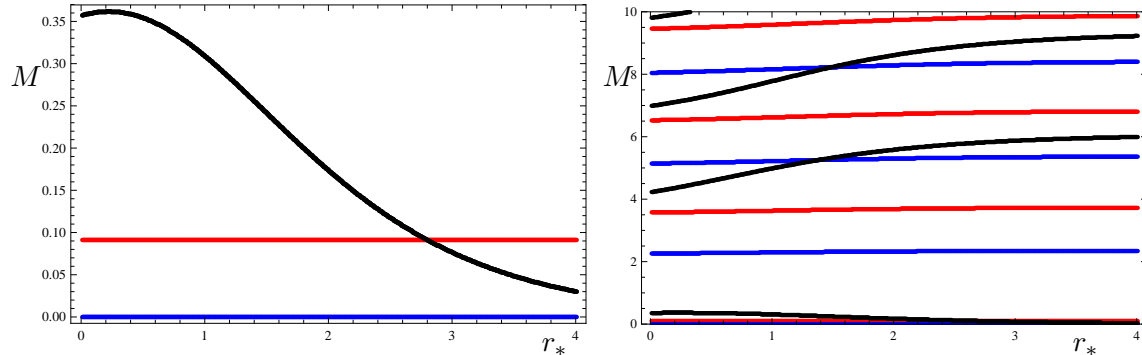


Figure 1: Left panel: the mass (in TeV) of the lightest states as a function of r_* , for the choices $r_1 = 0$, $\Delta = 1$ and $\Phi_I = 1$. Right panel: the mass spectrum of vectorial, axial and scalar states, for the same values of the parameters. In blue are the vector resonances, in red the axial-vector states, and in black the scalars.

enhance the coupling (and decay) to two photons. Similar arguments hold for the production mechanism, which results from competing suppression and enhancement mechanisms. Comparing this composite dilaton to the standard model Higgs has to be done with some caution, both from the theoretical point of view (a more precise and complete study of such couplings in this specific model would be necessary), and from the experimental point of view (a very precise measurement of the coupling would be needed).

The second problem arises from the fact that, because we assumed that symmetry-breaking be due to the IR boundary conditions (as opposed to the VEV of a bulk scalar), the only way to evade the bounds is to take ε^2 to be small. But in practice this means that the putative dual field theory is probably not at large- N , since $\varepsilon \propto \sqrt{N}$ (via unknown coefficients). In turn, this means that the classical gravity approximation we are using is probably at the border of its intrinsic validity. From the experimental perspective, this also means that the most accessible new state, the techni- ρ meson, decays mostly into a final state with two W bosons, and that this decay is not suppressed, so that the resonance is going to be somewhat broad. These two facts together might make the experimental detection of the techni- ρ quite challenging.

The reader should be aware of the fact that both these problems could be addressed by refining the model, and performing a more detailed study of all the couplings, which is beyond the purposes of this paper.

8 Conclusions.

The main physics message we want to convey, and that our paper exemplifies, is remarkably simple. Neither the discovery nor the exclusion of a light scalar (Higgs) particle provide, by themselves, sufficient evidence allowing to disentangle weakly-coupled and strongly-coupled origin of electro-weak symmetry breaking.

In the case where the Higgs particle is not found, it is not difficult to build models in which the couplings of such a state are suppressed, by assuming that several light scalars are present and that they mix non-trivially (an extreme example is for instance shown in [35]). On the other hand, the example in this paper illustrates how nowadays it is not difficult, thanks to gauge/gravity dualities, to build models in which a light dilaton (the couplings of which resemble the ones of the Higgs particle) is present in the spectrum,

but EWSB is triggered by a strongly-coupled interaction.

In order to distinguish the two scenarios, a thorough exploration of the mass range up to several TeV is necessary, and the LHC program has the potential to do so.

Note added: while in the process of finalizing this paper, the LHC collaborations announced an update on the results of the searches for the SM Higgs particle [36]. The results somewhat restrict the low-mass window allowed by the data, where a low-significance hint of a positive signal is present. These results leave our conclusions unchanged.

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